A common source for neutrino and sparticle masses[1]

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Abstract. We discuss supersymmetric scenarios in which neutrino masses arise from effective $d = 6$ operators in the Kähler potential (including SUSY-breaking insertions). Simple explicit realizations of those Kähler operators are presented in the context of the type II seesaw. An appealing scenario emerges upon identifying the seesaw mediators with SUSY-breaking messengers.

The smallness of neutrino masses can be explained by the conventional seesaw mechanism in which $m_\nu \sim v^2/M$, where v is the electroweak scale and $M \gg v$ is a heavy mass. It is also conceivable that neutrino masses are suppressed by a higher power of the heavy scale M, like $m_\nu \sim mv^2/M^2$, where $m \ll M$ is another mass parameter. In a SUSY framework this mass behaviour may stem from either superpotential or Kähler $d = 6$, $\Delta L = 2$ effective operators. Two examples of the latter have been proposed in [2], namely $\int d^4\theta\, (H_1^\dagger L)(H_2 L)/M^2$ and $\int d^4\theta\, (H_1^\dagger L)^2/M^2$, which imply $m \sim \mu$ (the superpotential Higgsino mass parameter). We point out the importance of including SUSY-breaking insertions to this approach and find new contributions to neutrino masses. We also present the simplest explicit realization of the Kähler operator $(H_1^\dagger L)^2$, including SUSY-breaking effects, and discuss a predictive scenario where the heavy seesaw mediators are also messengers of SUSY breaking.

1. Neutrino masses from Kähler operators and broken SUSY

Let us focus on the $\Delta L = 2$ effective operator $(H_1^\dagger L)^2/M^2$. In general, this will appear along with similar operators containing insertions of the type X/M_S, $X^\dagger/M_S$, $XX^\dagger/M_S^2$, where $X = \theta^2 F_X$ is a SUSY-breaking spurion superfield and M_S is the SUSY-breaking mediation scale, which can be either larger or smaller than M. Hence, the relevant $\Delta L = 2$ effective lagrangian can be parametrized as

$$\mathcal{L}_{\rm eff} = \int d^4\theta \frac{1}{2M^2}\left(\boldsymbol{\kappa} + \theta^2 \mathbf{B}_\kappa + \bar\theta^2 \tilde{\mathbf{B}}_\kappa + \theta^2\bar\theta^2 \mathbf{C}_\kappa\right)_{ij} (H_1^\dagger L_i)(H_1^\dagger L_j) \ + {\rm h.c.}\,, \qquad (1)$$

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where $\boldsymbol{\kappa}$ is dimensionless, $\mathbf{B}_\kappa$, $\tilde{\mathbf{B}}_\kappa$ and $\mathbf{C}_\kappa$ are dimensionful SUSY-breaking parameters and $i, j = e, \mu, \tau$ are flavour indices. The magnitude and flavour structure of all those parameters depend on the underlying physics which generates the effective operators.

The scale dependence of the above quantities is governed by their renormalization group equations (RGEs), which can be derived using the general expression of the one-loop corrected Kähler potential obtained in [3]. The RGE for $\boldsymbol{\kappa}$ can be found in [2], while the ones for the remaining operator coefficients have been obtained in [1]. For instance, the RGE for $\tilde{\mathbf{B}}_\kappa$ is:

$$
\begin{aligned}
8\pi^2 \frac{d\tilde{\mathbf{B}}_\kappa}{dt} &= \left[g^2 + g'^2 + \mathrm{Tr}(\mathbf{Y}_e^\dagger \mathbf{Y}_e + 3\mathbf{Y}_d^\dagger \mathbf{Y}_d)\right] \tilde{\mathbf{B}}_\kappa - \frac{1}{2}\left[\tilde{\mathbf{B}}_\kappa \mathbf{Y}_e^\dagger \mathbf{Y}_e + (\mathbf{Y}_e^\dagger \mathbf{Y}_e)^T \tilde{\mathbf{B}}_\kappa\right] \\
&\quad + \left[g^2 M_2^* + g'^2 M_1^* - 2\,\mathrm{Tr}(\mathbf{A}_e^\dagger \mathbf{Y}_e + 3\mathbf{A}_d^\dagger \mathbf{Y}_d)\right] \boldsymbol{\kappa} + \boldsymbol{\kappa}\, \mathbf{A}_e^\dagger \mathbf{Y}_e + (\mathbf{A}_e^\dagger \mathbf{Y}_e)^T \boldsymbol{\kappa}\,. \qquad (2)
\end{aligned}
$$

All four operators of eq. (1) contribute to neutrino masses. Those with coefficients $\boldsymbol{\kappa}$ and $\tilde{\mathbf{B}}_\kappa$ contribute to $\mathbf{m}_\nu$ at the tree-level, in such a way that $\mathbf{m}_\nu \simeq \mathbf{m}_\nu^{(\kappa)} + \mathbf{m}_\nu^{(\tilde{B}_\kappa)}$. The $\boldsymbol{\kappa}$-operator leads to a lagrangian term of the form $(F_{H_1}^\dagger L)(H_1^\dagger L)$, which contributes to neutrino masses as [2]

$$
\mathbf{m}_\nu^{(\kappa)} = 2\,\boldsymbol{\kappa}\,\mu \frac{v^2}{M^2} \sin\beta\cos\beta\,. \qquad (3)
$$

The $\tilde{\mathbf{B}}_\kappa$-operator gives a lagrangian term of the form $(H_1^\dagger L)^2$, which induces [1]

$$
\mathbf{m}_\nu^{(\tilde{B}_\kappa)} = \tilde{\mathbf{B}}_\kappa \frac{v^2}{M^2} \cos^2\beta\,. \qquad (4)
$$

This novel contribution to $\mathbf{m}_\nu$ can dominate over $\mathbf{m}_\nu^{(\kappa)}$. The remaining operators of eq. (1), with coefficients $\mathbf{B}_\kappa$ and $\mathbf{C}_\kappa$, contribute to neutrino masses through finite one-loop diagrams involving gauginos and sleptons. Considering the soft mass matrix of 'left-handed' sleptons $\tilde{L}$ as $\mathbf{m}_{\tilde{L}}^2 = \tilde{m}_L^2(\mathbb{1} + \boldsymbol{\Delta}_L)$ ($\tilde{m}_L^2$ sets the overall mass scale and the dimensionless matrix $\boldsymbol{\Delta}_L$ accounts for the flavour dependence), we get the following contributions to $\mathbf{m}_\nu$ (at first order in $\boldsymbol{\Delta}_L$)

$$
\delta_{B_\kappa}\mathbf{m}_\nu \simeq \tfrac{1}{32\pi^2}\left[-\left(\tfrac{g^2}{M_2} f_{L2} + \tfrac{g'^2}{M_1} f_{L1}\right)\mathbf{B}_\kappa + \left(\tfrac{g^2}{M_2} h_{L2} + \tfrac{g'^2}{M_1} h_{L1}\right)(\mathbf{B}_\kappa \boldsymbol{\Delta}_L + \boldsymbol{\Delta}_L^T \mathbf{B}_\kappa)\right] 2\mu \tfrac{v^2}{M^2} \sin\beta\cos\beta\,, \qquad (5)
$$

$$
\delta_{C_\kappa}\mathbf{m}_\nu \simeq \tfrac{1}{32\pi^2}\left[-\left(\tfrac{g^2}{M_2} f_{L2} + \tfrac{g'^2}{M_1} f_{L1}\right)\mathbf{C}_\kappa + \left(\tfrac{g^2}{M_2} h_{L2} + \tfrac{g'^2}{M_1} h_{L1}\right)(\mathbf{C}_\kappa \boldsymbol{\Delta}_L + \boldsymbol{\Delta}_L^T \mathbf{C}_\kappa)\right] \tfrac{v^2}{M^2} \cos^2\beta\,, \qquad (6)
$$

where $f_{La} = f(\tilde{m}_L^2/|M_a|^2)$, $h_{La} = h(\tilde{m}_L^2/|M_a|^2)$, $f(x) = (x - 1 - \log x)/(x-1)^2$ and $h(x) = (x^2 - 1 - 2x\log x)/(x-1)^3$. Both the flavour structure and the size of $\delta_{B_\kappa}\mathbf{m}_\nu$, $\delta_{C_\kappa}\mathbf{m}_\nu$ are model dependent.

2. Type II seesaw realizations

Among the three variants of the seesaw mechanism (which generate the familiar $d = 5$ superpotential operator at the tree level), the type II is the natural one in which the above $d = 6$ operators emerge. In fact, the tree-level exchange of type I or type III mediators leads to $\Delta L = 0$ Kähler operators of the form $|H_2 L|^2$, whereas the type II mediators induce both $\Delta L = 0$ and $\Delta L = 2$ operators. The type II seesaw mechanism is realized through the exchange of $SU(2)_W$ triplet states T and $\bar{T}$ in a vector-like $SU(2)_W \times U(1)_Y$ representation, $T \sim (3, 1)$, $\bar{T} \sim (3, -1)$. The relevant superpotential is: $W \supset \frac{1}{\sqrt{2}}\mathbf{Y}_T^{ij} L_i T L_j + \frac{1}{\sqrt{2}}\lambda_1 H_1 T H_1 + \frac{1}{\sqrt{2}}\lambda_2 H_2 \bar{T} H_2 + M_T T\bar{T}$, where $\mathbf{Y}_T^{ij}$ is a 3×3 symmetric matrix, $\lambda_{1,2}$ are dimensionless couplings and M_T is the (SUSY) triplet mass. By integrating out the triplets one obtains $\Delta L = 2$ effective operators of dimension

$d=5$, *i.e.* $W_{\rm eff}\supset \frac{\lambda_2}{2M_T}\mathbf{Y}_T^{ij}(L_iH_2)(L_jH_2)$, and $d=6$, *i.e.* $K_{\rm eff}\supset \frac{\lambda_1^*}{2|M_T|^2}\mathbf{Y}_T^{ij}(H_1^\dagger L_i)(H_1^\dagger L_j)+h.c.$, both of which can generate neutrino masses. The former operator is usually leading, but we assume it to be strongly suppressed (absent) by a very small (vanishing) value of λ_2, which can be justified by symmetry arguments. In this case the leading $\Delta L=2$ operator is the Kähler one, which can be matched to the SUSY part of eq. (1) through the identification: $\boldsymbol{\kappa}=\lambda_1^*\mathbf{Y}_T$ and $M^2=|M_T|^2$. The resulting contribution to neutrino masses is the tree-level term $\mathbf{m}_\nu^{(\kappa)}$ of eq. (3).

One can now ask how the SUSY-breaking operators of eq. (1) arise in the type II seesaw framework. The answer to this question depends on the ordering of the SUSY-breaking mediation scale M_S and the triplet mass M_T: $M_S<M_T$; $M_S>M_T$ or $M_S=M_T$.

• $M_S<M_T$. If SUSY breaking mediation occurs at $M_S<M_T$ via a messenger sector coupled to the MSSM through gauge interactions only (pure gauge mediation), then SUSY-breaking gaugino and (flavour blind) sfermion masses arise at M_S through loop diagrams, while trilinear terms are driven below M_S by gaugino mass terms in the RGEs. As for our $\Delta L=2$ operators of eq. (1), only the SUSY one with coefficient $\boldsymbol{\kappa}$ is present above M_S (it is generated at M_T), whereas the SUSY-breaking ones receive finite two-loop contributions at M_S, and logarithmic ones below M_S through RGEs. The dominant source of neutrino masses is, generically, the SUSY contribution $\mathbf{m}_\nu^{(\kappa)}$ of eq. (3).

• $M_S>M_T$. Suppose that SUSY-breaking terms are generated at $M_S>M_T$ through, *e.g.*, gravity or gauge mediation. This means that all the MSSM and triplet fields have SUSY-breaking mass parameters at M_T. In this case, the tree-level decoupling of the triplets generates all the $\Delta L=2$ effective operators of eq. (1), with coefficients $\boldsymbol{\kappa}=\lambda_1^*\mathbf{Y}_T$, $\mathbf{B}_\kappa=\lambda_1^*(\mathbf{Y}_TB_T-\mathbf{A}_T)$, $\tilde{\mathbf{B}}_\kappa=(\lambda_1^*B_T^*-A_1^*)\mathbf{Y}_T$, $\mathbf{C}_\kappa=(\lambda_1^*B_T^*-A_1^*)(\mathbf{Y}_TB_T-\mathbf{A}_T)-\lambda_1^*\mathbf{Y}_Tm_{\tilde T}^2$.

3. $M_S=M_T$: Seesaw mediators as SUSY-breaking messengers

The $M_S=M_T$ scenario is obtained by identifying T and $\bar T$ as being SUSY-breaking mediators. They are embedded in a messenger sector which (in order to generate the gluino mass) should also include coloured fields. We also require that perturbative unification of gauge couplings be preserved and that all messenger masses be of the same order. This implies that the messenger sector should have a common total Dynkin index N for each subgroup of $SU(3)_C\times SU(2)_W\times U(1)_Y$. Since the $T+\bar T$ pair has $SU(2)$ index $N_2=4$, we are constrained to $N\geq 4$. The minimal case ($N=4$) can be built by adding a pair of $SU(3)_C$ triplets $(3,1,-1/3)+(\bar 3,1,+1/3)$ and an $SU(3)_C$ adjoint $(8,1,0)$ to T and $\bar T$.

The MSSM SUSY-breaking parameters are generated at the quantum level by a messenger sector of the type described above, coupled to the MSSM fields through both gauge and Yukawa interactions [4]. At the one-loop level, the gaugino masses M_a, the Higgs B-term B_H and the trilinear terms $\mathbf{A}_x$ are: $M_a=-\frac{NB_T}{16\pi^2}g_a^2$, $B_H=\frac{3B_T}{16\pi^2}|\lambda_1|^2$ $\mathbf{A}_e=\frac{3B_T}{16\pi^2}\mathbf{Y}_e(\mathbf{Y}_T^\dagger\mathbf{Y}_T+|\lambda_1|^2)$, $\mathbf{A}_d=\frac{3B_T}{16\pi^2}\mathbf{Y}_d|\lambda_1|^2$ and $\mathbf{A}_u=0$, where $g_1^2=(5/3)g'^2$ and $g_2^2=g^2$. Non-vanishing $\mathcal{O}(B_T^2)$ contributions for the squared scalar masses arise at the two-loop level. For sleptons these are:

$$\begin{aligned}\mathbf{m}^2_{\tilde L} &= \left(\frac{|B_T|}{16\pi^2}\right)^2\Big[N\left(\frac{3}{10}g_1^4+\frac{3}{2}g_2^4\right)-\left(\frac{27}{5}g_1^2+21g_2^2\right)\mathbf{Y}_T^\dagger\mathbf{Y}_T+3|\lambda_1|^2(\mathbf{Y}_T^\dagger\mathbf{Y}_T-\mathbf{Y}_e^\dagger\mathbf{Y}_e)\\ &\qquad +3\,\mathbf{Y}_T^\dagger(\mathbf{Y}_e^\dagger\mathbf{Y}_e)^T\mathbf{Y}_T+18\,(\mathbf{Y}_T^\dagger\mathbf{Y}_T)^2+3\,\mathbf{Y}_T^\dagger\mathbf{Y}_T\mathrm{Tr}(\mathbf{Y}_T^\dagger\mathbf{Y}_T)\Big]\,, \qquad (7)\\ \mathbf{m}^2_{\tilde e^c} &= \left(\frac{|B_T|}{16\pi^2}\right)^2\Big[N\left(\frac{6}{5}g_1^4\right)-6\,\mathbf{Y}_e(\mathbf{Y}_T^\dagger\mathbf{Y}_T+|\lambda_1|^2)\mathbf{Y}_e^\dagger\Big]\,. \qquad (8)\end{aligned}$$

Notice that the flavour structures of $\mathbf{A}_e$, $\mathbf{m}^2_{\tilde L}$ and $\mathbf{m}^2_{\tilde e^c}$ are controlled by $\mathbf{Y}_T$ and $\mathbf{Y}_e$, which in turn are determined by the low-energy lepton masses and mixing angles. Such minimal LFV

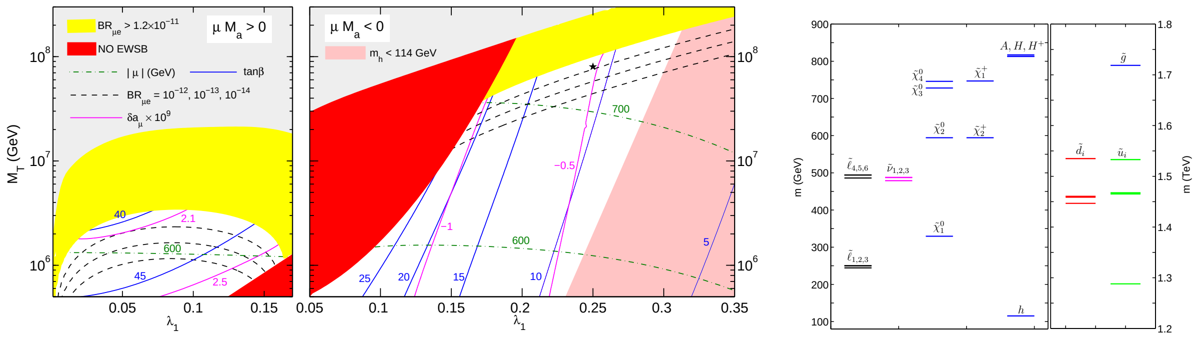

Figure 1. Plots of the $N = 4$ model for $B_T = 60$ TeV and normally ordered neutrino spectrum with $0 = m_1 < m_2 \ll m_3$ and $s_{13} = 0$. *Left panels*: The (λ_1, M_T) parameter space for $\mu M_a > 0$ (left) and $\mu M_a < 0$ (right). The white region is the allowed one. The dashed lines correspond to $\mathrm{BR}(\mu \to e\gamma) = 10^{-12}, 10^{-13}, 10^{-14}$ (from top to bottom.). *Right panel*: Sparticle and Higgs spectrum for $M_T = 8 \times 10^7$ GeV and $\lambda_1 = 0.25$ (point $\star$ of the parameter space).

properties are a characteristic feature of the SUSY type II seesaw [4, 5]. Our scenario is a flavoured variant of gauge mediation and possesses the property of minimal flavour violation in both the quark and lepton sectors.

Upon decoupling the triplets, the MSSM SUSY-breaking masses are generated through finite radiative effects, while the $\Delta L = 2$ SUSY-breaking parameters $\mathbf{B}_\kappa$, $\tilde{\mathbf{B}}_\kappa$ and $\mathbf{C}_\kappa$ arise at the tree level. The latter have a very simple form: $\mathbf{B}_\kappa = B_T \boldsymbol{\kappa}$, $\tilde{\mathbf{B}}_\kappa = B_T^* \boldsymbol{\kappa}$ and $\mathbf{C}_\kappa = |B_T|^2 \boldsymbol{\kappa}$, where $\boldsymbol{\kappa} = \lambda_1^* \mathbf{Y}_T$. Eqs. (3) and (4) imply that the tree-level contribution to $\mathbf{m}_\nu$ is:

$$\mathbf{m}_\nu = \boldsymbol{\kappa} \,(B_T^* + 2\,\mu \tan\beta) \cos^2\!\beta \,\frac{v^2}{|M_T|^2}\,. \tag{9}$$

Notice that the SUSY-breaking parameter B_T acts as a common source for both sparticle and neutrino masses.

This scenario has a small number of free parameters, namely M_T, B_T, λ_1 and the messenger index N. Once these are fixed, the remaining parameters $\mathbf{Y}_T$, $\tan\beta$ and μ are determined by the low-energy neutrino data and by requiring proper EWSB. Concerning neutrino data, we recall that the neutrino mass matrix $\mathbf{m}_\nu$ is related to the low-energy observables as $\mathbf{m}_\nu = \mathbf{U}^* \mathbf{m}_\nu^D \mathbf{U}^\dagger$, where $\mathbf{m}_\nu^D = \mathrm{diag}(m_1, m_2, m_3)$, m_a are the neutrino masses and $\mathbf{U}$ is the lepton mixing matrix. Some representative numerical results for the $N = 4$ model are shown in Fig. 1, for $B_T = 60$ TeV and a normally ordered neutrino spectrum, with $0 = m_1 < m_2 \ll m_3$ and $s_{13} = 0$. In the left part of Fig. 1 we show two plots of the (λ_1, M_T) parameter space, including contours of $\tan\beta$ and μ (extracted by imposing EWSB). The left (right) panel corresponds to solutions of the EWSB conditions with $\mu M_a > 0$ ($\mu M_a < 0$). The main phenomenological constraints come from the LFV decay $\mu \to e\gamma$ and the lightest Higgs mass. In the right panel of Fig. 1 we show the sparticle and Higgs spectrum for $M_T = 8 \times 10^7$ GeV and $\lambda_1 = 0.25$ (which corresponds to $\tan\beta \simeq 11$), again for $B_T = 60$ TeV. The Higgs sector is in the decoupling regime, since the states A, H and H^+ are much heavier than h. Gluino and squarks are the heaviest sparticles and the lightest of them is $\tilde{t}_1$ (which is mainly $\tilde{t}_R$). In the electroweak sector, the heaviest chargino and neutralinos ($\tilde{\chi}_2^+$, $\tilde{\chi}_{3,4}^0$) are mainly Higgsino-like, while $\tilde{\chi}_1^+$ and $\tilde{\chi}_2^0$ are mostly Wino-like. The (mainly left-handed) sleptons $\tilde{\ell}_{4,5,6}$ and the sneutrinos $\tilde{\nu}_{1,2,3}$ are somewhat lighter than those states, and the Bino-like neutralino $\tilde{\chi}_1^0$ is even lighter. Finally, the lightest MSSM sparticles are the (mainly right-handed) sleptons $\tilde{\ell}_{1,2,3}$, as generically occurs in gauge mediated models with

messenger index $N > 1$ and not too large mediation scale. The slepton $\tilde{\ell}_1$ (which is mainly $\tilde{\tau}_R$) is the next-to-lightest SUSY particle (NLSP), while the gravitino $\tilde{G}$ is the lightest one (LSP).

The scenario described above can be tested at current and future colliders. For instance, if a pp collision at the LHC produces two squarks, each of them can decay through well known chains, such as $\tilde{q}_R \to q\tilde{\chi}^0_1 \to q\tau\tilde{\ell}_1$, or $\tilde{q}_L \to q\tilde{\chi}^0_2 \to q\ell\tilde{\ell} \to q\ell^+\ell^-\tilde{\chi}^0_1 \to q\ell^+\ell^-\tau\tilde{\ell}_1$, or similar ones with charginos and/or sneutrinos (and neutrinos). Hence, in general, the final state of such a collision contains SM particles and two NLSPs $\tilde{\ell}_1$, which eventually decay to $\tau\tilde{G}$ with rate $\Gamma = m^5_{\tilde{\ell}_1}/(16\pi F^2)$. The latter decays can occur either promptly, or at a displaced vertex, or even outside the main detector.

Since LFV is intrinsically present in our framework, LFV processes are a crucial tool to discriminate our model from pure gauge mediation ones. For moderate $\tan\beta$, the leading LFV structure $\mathbf{Y}_T^\dagger\mathbf{Y}_T$, which appears in $\mathbf{m}^2_{\tilde{L}}$ (and $\mathbf{A}_e$), can be related to the neutrino parameters as

$$(\mathbf{m}^2_{\tilde{L}})_{ij} \propto B_T^2(\mathbf{Y}_T^\dagger\mathbf{Y}_T)_{ij} \propto \left(\frac{M_T^2\tan^2\beta}{\lambda_1}\right)^2 \left[\mathbf{V}(\mathbf{m}^D_\nu)^2\mathbf{V}^\dagger\right]_{ij} \propto \tan^5\beta\, M_T^4 \left[\mathbf{V}(\mathbf{m}^D_\nu)^2\mathbf{V}^\dagger\right]_{ij} . \quad (10)$$

LFV signals can therefore appear either at high-energy colliders or in low-energy processes. Concerning the former possibility, LFV could show up at the LHC in, *e.g.*, neutralino decays, such as $\tilde{\chi}^0_2 \to \ell_i^\pm\tilde{\ell}_a^\mp \to \ell_i^\pm\ell_j^\mp\tilde{\chi}^0_1$ with $i \neq j$, followed by the flavour-conserving decay $\tilde{\chi}^0_1 \to \tau^+\tau^-\tilde{G}$ (or $\tilde{\chi}^0_1 \to \tilde{\ell}_1^\pm\tau^\mp$ if the NLSP is long-lived).

As for low-energy LFV processes, let us focus on the radiative decays $\ell_i \to \ell_j\gamma$. By using eq. (10), we can infer that

$$\mathrm{BR}(\ell_i \to \ell_j\gamma) \propto \frac{|(\mathbf{m}^2_{\tilde{L}})_{ij}|^2}{\tilde{m}^8}\tan^2\beta \propto \left(\frac{M_T}{B_T}\right)^8 (\tan\beta)^{12} \left|\left[\mathbf{V}(\mathbf{m}^D_\nu)^2\mathbf{V}^\dagger\right]_{ij}\right|^2 , \quad (11)$$

where the flavour-dependence is determined by the low-energy neutrino parameters only [5, 4]. If we take ratios of BRs, we have that, for a normal neutrino mass spectrum, $\mathrm{BR}(\tau \to \mu\gamma)/\mathrm{BR}(\mu \to e\gamma) \simeq 400\,(2)\,[3]$ and $\mathrm{BR}(\tau \to e\gamma)/\mathrm{BR}(\mu \to e\gamma) \simeq 0.2\,(0.1)\,[0.3]$ if $s_{13} = 0\,(s_{13} = 0.2, \delta = 0)\,[s_{13} = 0.2, \delta = \pi]$. These approximate results hold for small or moderate $\tan\beta$. Special features may emerge for large $\tan\beta$ [6].

In the left panel of Fig. 2 the three $\mathrm{BR}(\ell_i \to \ell_j\gamma)$ are shown as a function of $\tan\beta$, taking either $B_T = 60\,\mathrm{TeV}$ (solid lines) or $B_T = 100\,\mathrm{TeV}$ (dashed lines), for $M_T = 8 \times 10^7\,\mathrm{GeV}$, normally ordered neutrino spectrum and $s_{13} = 0$. In both examples $\mathrm{BR}(\mu \to e\gamma)$ can be tested at the MEG experiment for suitable ranges of $\tan\beta$. If $\mathrm{BR}(\mu \to e\gamma)$ is close to its present bound, $\mathrm{BR}(\tau \to \mu\gamma)$ is above 10^{-9}, within the reach of future Super Flavour Factories.

The double panel on the right of Fig. 2 illustrates the dependence of the BRs on the least known neutrino parameters, namely s_{13} and δ, for $B_T = 60\,\mathrm{TeV}$, $M_T = 8 \times 10^7\,\mathrm{GeV}$ and two values of $\tan\beta$. Regarding $\mathrm{BR}(\mu \to e\gamma)$ and $\mathrm{BR}(\tau \to e\gamma)$, the solid (dashed) curves correspond to $\delta = 0\,(\pi)$, while the region between such curves is spanned by intermediate values of δ. The dependence of $\mathrm{BR}(\tau \to \mu\gamma)$ on s_{13} and δ is negligible. The first subpanel corresponds to a scenario which could be tested very soon at MEG through the search of $\mu \to e\gamma$, if $s_{13} \ll 0.01$. Notice that $\mathrm{BR}(\tau \to \mu\gamma)$ in this example is around 4×10^{-9}, within the reach of future Super Flavour Factories, while $\tau \to e\gamma$ would be unobservable because $\mathrm{BR}(\tau \to e\gamma) \sim 10^{-12}$. For $s_{13} \sim 0.01$, $\mathrm{BR}(\mu \to e\gamma)$ and $\mathrm{BR}(\tau \to e\gamma)$ can be either enhanced or suppressed since, depending on the value of δ, a cancellation can occur in the LFV quantity $\left[\mathbf{V}(\mathbf{m}^D_\nu)^2\mathbf{V}^\dagger\right]_{ij}$ [4, 6]. The cancellation takes place in $\mathrm{BR}(\mu \to e\gamma)$ $[\mathrm{BR}(\tau \to e\gamma)]$ for $\delta = \pi\,(0)$ in the case of normal ordering, while the opposite occurs for inverted ordering. If $\mathrm{BR}(\mu \to e\gamma)$ is suppressed by that cancellation mechanism, only $\tau \to \mu\gamma$ can be observed. In such a case, we can even obtain

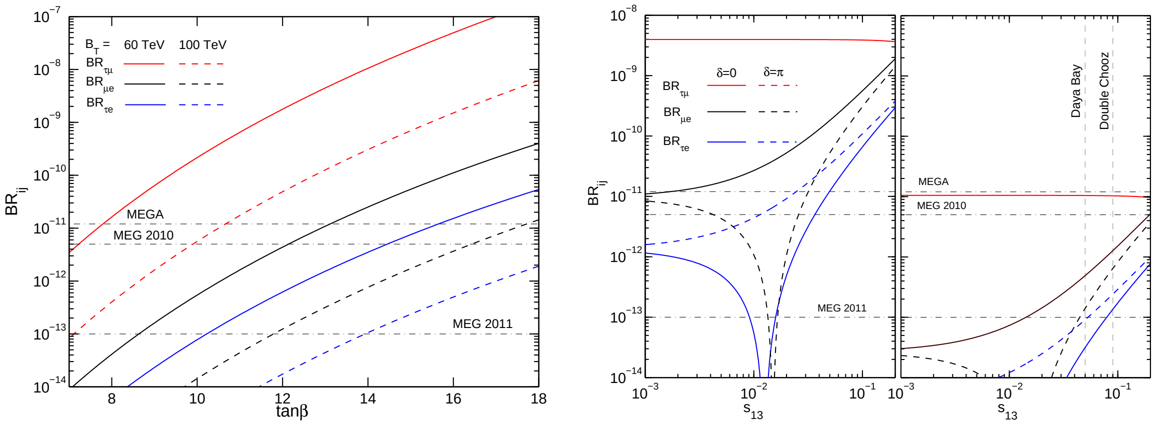

Figure 2. Plots of $\mathrm{BR}(\ell_i \to \ell_j\gamma)$ for $M_T = 8 \times 10^7\,\mathrm{GeV}$. *Left panel*: The BRs as a function of $\tan\beta$ for $B_T = 60\,\mathrm{TeV}$ (solid lines) or $B_T = 100\,\mathrm{TeV}$ (dashed lines), with $s_{13} = 0$. *Right panel*: The BRs as a function of s_{13} for $B_T = 60\,\mathrm{TeV}$, with $\tan\beta = 13\,(8)$ in the first (second) subpanel. For $\mathrm{BR}(\mu \to e\gamma)$ and $\mathrm{BR}(\tau \to e\gamma)$, the solid (dashed) curves correspond to $\delta = 0\,(\pi)$ assuming a normally-ordered neutrino spectrum.

values of $\mathrm{BR}(\tau \to \mu\gamma)$ above 10^{-8} by slightly changing the model parameters. In the case of partial cancellations, $\mu \to e\gamma$ could be still probed by MEG for values of s_{13} up to about 0.03, which are in the potential reach of future Neutrino Factories. The second subpanel shows an alternative possibility, in which LFV τ decays are unobservable, whereas $\mathrm{BR}(\mu \to e\gamma)$ lies in the range $10^{-13} - 5 \times 10^{-12}$ if $0.05 \sim s_{13} < 0.2$. Those values of $\mathrm{BR}(\mu \to e\gamma)$ should be probed by MEG next year, while the indicated range of s_{13} is within the sensitivity of the present and incoming neutrino experiments. This example shows the importance of the interplay between LFV searches and neutrino oscillation experiments.

In conclusion, we have summarized the investigation of Ref. [1] on an interesting alternative neutrino mass mechanism, which relies on the $d = 6$, $\Delta L = 2$ Kähler operator $(H_1^\dagger L)^2/M^2$. We have presented both a general effective-theory description and explicit realizations in the context of the type II seesaw mechanism. If the seesaw mediators are also identified with SUSY-breaking messengers, strong correlations arise between neutrino parameters, sparticle and Higgs masses, as well as LFV processes.

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